

Neutrinoless double- β decay in TeV scale Left-Right symmetric models

Joydeep Chakraborty,^{1,*} H. Zeen Devi,^{1,†} Srubabati Goswami,^{1,‡} and Sudhanwa Patra^{2,§}

¹*Physical Research Laboratory, Ahmedabad-380009, India*

²*Institute of Physics, Bhubaneswar-, India*

Abstract

In this paper we study in detail the neutrinoless double beta decay in left-right symmetric models with right-handed gauge bosons at TeV scale which is within the presently accessible reach of colliders. We discuss the different diagrams that can contribute to this process and identify the dominant ones for the case where the right-handed neutrino is also at the TeV scale. We calculate the contribution to the effective mass governing neutrinoless double beta decay assuming type-I, and type-II dominance and discuss what are the changes in the effective mass due to the additional contributions. We also discuss the effect of the recent Daya-Bay and RENO measurements on $\sin^2 \theta_{13}$ on the effective mass in different scenarios.

*Electronic address: joydeep@prl.res.in

†Electronic address: zeen@prl.res.in

‡Electronic address: sruba@prl.res.in

§Electronic address: sudha.astro@gmail.com

I. INTRODUCTION

Neutrino oscillation experiments have steadfastly established that neutrinos are massive and mix between different flavours. The Standard Model (SM) does not accommodate any right-handed (RH) partner for the neutrino and hence they remain massless. Thus, existence of neutrino mass requires one to tread beyond the standard paradigm. A prevalent way to generate small neutrino mass is through the seesaw mechanism. This is due the Weinberg operator $\kappa LLHH$ [1] where L and H are respectively lepton and Higgs fields transforming as $SU(2)$ doublets; κ is the effective coupling which has inverse mass dimension. Such terms violate lepton number by 2 units and hence essentially predict neutrinos to be Majorana particles. Violation of lepton number can be manifested in neutrinoless double beta decay ($0\nu\beta\beta$) process :

$$(A, Z) \rightarrow (A, Z + 2) + 2e^-$$

which is therefore of prime importance as it can ascertain the nature of the neutrinos and hence the mechanism of neutrino mass generation. The best limit on the half life of this process at present is $T_{1/2} < 3 \times 10^{25}$ years coming from the Heidelberg-Moscow experiment using ^{76}Ge [2]. This translates to a bound on the effective neutrino mass [3]

$$m_{\text{eff}} \leq 0.21 - 0.53 \text{ eV} \tag{1}$$

where the range is due to the uncertainty in the nuclear matrix elements. There are many other upcoming experiments to improve this bound [4] and also to test the claim of the positive evidence for $0\nu\beta\beta$ by a part of the Heidelberg-Moscow Collaboration [5, 6].

If the only contribution to $0\nu\beta\beta$ is through exchange of light neutrinos then the effective neutrino mass is just the absolute value of the (ee) element of the low energy neutrino mass matrix in the flavour basis. However there can be other scenarios like models with heavy right-handed neutrinos [7–9], R-parity violating Supersymmetry [10–13], extra dimensional [14] scenarios etc. which can give rise to additional diagrams contributing to the neutrinoless double beta decay process. Of special importance are the scenarios where the scale of new physics is at $\sim \text{TeV}$ since then the additional contributions to $0\nu\beta\beta$ can be significant [15–17]. Such models have recently gained prominence since it these can be probed at the LHC and can also give rise to Lepton Flavour Violating processes[18]. It is noteworthy to mention at this point that lepton number violation can also be probed at LHC through the so called

golden process [19] producing same sign di-leptons in the final state and complimentary information from LHC and $0\nu\beta\beta$ may help in consolidating the nature of new physics and origin of neutrino mass.

In this paper we focus on the TeV scale left-right symmetric seesaw models for generation of neutrino mass and the implications for $0\nu\beta\beta$. Left-right (LR) symmetric models were motivated from the perspective of restoring parity symmetry at a high scale. This required placing the left- and right-handed fermions on the same footing by putting them as part of $SU(2)_L$ and $SU(2)_R$ doublets respectively [20]. The observation of parity violation in weak interaction indicates that the LR symmetry is broken at a low scale and in the Minimal LR symmetric model this is achieved by triplet scalar fields. Such a framework naturally generates neutrino masses through type-I seesaw due to the right-handed neutrinos [21] and type-II seesaw via the triplet scalars [22].

Because $SU(2)_R$ is a gauge symmetry the LR symmetric models contain right-handed charged currents mediated by W_R boson. If the right-handed W_R masses are around TeV then it is possible to access them at LHC. This has motivated several studies recently with the right-handed W_R at the TeV scale [23–25]. In such a scenario, one can have contributions to $0\nu\beta\beta$ from both left-handed (LH) and right-handed currents via exchange of light and heavy neutrinos respectively. Additionally there can be diagrams mediated by the charged Higgs fields. Our work scrutinizes in detail the relative contributions of various diagrams for TeV scale LR symmetric models. We consider seesaw scenarios with type-I, and type-II dominance and discuss the behaviour of the effective mass in different limits. The implications for type-II seesaw dominance in the context of LR symmetric models have been considered in [16]. In our work we present the analytic expressions of the effective mass in different limits and discuss their dependence on the neutrino oscillation parameters. In particular we incorporate the recent results on θ_{13} measurement from Daya-Bay [26] and RENO [27] experiment and discuss the consequences. We also study the contribution of the triplet Higgs mediated diagrams and discuss for which scenarios it can contribute.

The plan of the paper is as follows. In the next section we discuss type-I, and type-II seesaw in LR symmetric model. In section 3 we consider neutrinoless double beta decay due to light neutrinos. In section 4 we examine the other additional contributions to $0\nu\beta\beta$ within the LR model in both type I, and type II dominance. We also summarize briefly the possible impact of the Higgs triplet contribution at the end of this section. In the Appendix

we present the Feynman diagrams and estimated different contributions. We conclude with an overview of our study.

II. NEUTRINO MASS IN LR SYMMETRIC MODEL

In left-right symmetric models, the standard model gauge group is augmented to include a right-handed $SU(2)$ counterpart enlarging the group to $SU(3)_c \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$. Thus along with the left-handed doublet fermions we also have replicas that transform as doublets under $SU(2)_R$. The $SU(2)_R$ is related to the $SU(2)_L$, by a discrete symmetry which is the anthem of the left-right symmetric model. We restrict ourselves within the so-called minimal form of the model that includes a bidoublet(Φ), triplet($\Delta_{L/R}$)-scalars, right-handed neutrino(N_R), extra gauge bosons(W_R^\pm, Z_2) along with the SM particles. The assigned quantum numbers of these extra scalars under $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$ read as:

$$\Phi \equiv (2, 2, 0), \Delta_L \equiv (3, 1, 1), \Delta_R \equiv (1, 3, 1). \quad (2)$$

The Lagrangian giving the neutral fermion mass terms is

$$\mathcal{L}_Y = f_L l_L^T C i \sigma_2 \Delta_L l_L + f_R l_R^T C i \sigma_2 \Delta_R l_R + \bar{l}_R (y_D \phi + y_L \tilde{\Phi}) l_L + h.c., \quad (3)$$

where $l_L(l_R)$ denotes the left(right)- handed fermion doublets and $\tilde{\Phi} = \sigma_2 \Phi^* \sigma_2$. The bi-doublet Higgs acquires vacuum expectation values (vev) as:

$$\langle \Phi \rangle = \begin{pmatrix} v & 0 \\ 0 & v' \end{pmatrix}. \quad (4)$$

The charge lepton masses are given as:

$$m_l = y_D v' + y_L v \quad (5)$$

$$m_D = y_L v' + y_D v. \quad (6)$$

Here, we consider $y_D v \gg y_L v'$, and $y_L v \gg y_D v'$ so that the Yukawa coupling matrix responsible for charged fermion masses is y_L while for the neutrino masses it is y_D .

Once the Higgs fields develop vev the neutral fermion mass matrix becomes

$$M_\nu \equiv \begin{pmatrix} f_L v_L & y_D v \\ y_D^T v & f_R v_R \end{pmatrix}, \quad (7)$$

where $\langle \Delta_L \rangle = v_L$, $\langle \Delta_R \rangle = v_R$. After block diagonalizing M_ν in the seesaw approximation ($f_R v_R \gg y_D v$) we get

$$(m_\nu^{light})_{3 \times 3} = f_L v_L + \frac{v^2}{v_R} y_D^T f_R^{-1} y_D, \quad (8)$$

$$(m_R^{heavy})_{3 \times 3} = f_R v_R, \quad (9)$$

where, the first term in Eq.(8) is due to type-II seesaw whereas the second term corresponds to type-I seesaw mediated by right-handed neutrinos.

In this model the Yukawa couplings f_L and f_R are not independent but related to each other. We can have either $f_L = f_R$, and $U_L = U_R$ or $f_L = f_R^*$, and $U_L = U_R^*$ as an artifact of the LR-symmetry. This constrains the extra freedom in the Yukawa sector and reduces the number of free parameters.

III. NEUTRINOLESS DOUBLE- β DECAY IN THREE GENERATION PICTURE

In the standard three generation picture the time period for neutrinoless double beta decay is given as,

$$\frac{\Gamma_{0\nu\beta\beta}}{\ln 2} = G \left| \frac{\mathcal{M}_\nu}{m_e} \right|^2 |m_\nu^{ee}|^2, \quad (10)$$

where G contains the phase space factors, m_e is the electron mass, \mathcal{M}_ν is the nuclear matrix element.

$$|m_\nu^{ee}| = |U_{ei}^2 m_i|, \quad (11)$$

is the effective neutrino mass that appear in the expression for time period for neutrinoless double beta decay.

The unitary matrix U is the so called PMNS mixing matrix which coincides with the neutrino mixing matrix in the basis where charged lepton mass matrix is diagonal. The standard parametrization for this is

$$U = \begin{pmatrix} c_{12} c_{13} & s_{12} c_{13} & s_{13} e^{-i\delta} \\ -c_{23} s_{12} - s_{23} s_{13} c_{12} e^{i\delta} & c_{23} c_{12} - s_{23} s_{13} s_{12} e^{i\delta} & s_{23} c_{13} \\ s_{23} s_{12} - c_{23} s_{13} c_{12} e^{i\delta} & -s_{23} c_{12} - c_{23} s_{13} s_{12} e^{i\delta} & c_{23} c_{13} \end{pmatrix} P. \quad (12)$$

The abbreviations used above are $c_{ij} = \cos \theta_{ij}$, $s_{ij} = \sin \theta_{ij}$, δ is the Dirac CP phase while the phase matrix $P = \text{diag}(1, e^{i\alpha_2}, e^{i(\alpha_3+\delta)})$ contains the Majorana phases α_2 and α_3 . The 3

parameter	best-fit	3σ
$\Delta m_{\text{sol}}^2 [10^{-5} \text{eV}^2]$	7.58	6.99-8.18
$ \Delta m_{\text{atm}}^2 [10^{-3} \text{eV}^2]$	2.35	2.06-2.67
$\sin^2 \theta_{12}$	0.306	0.259-0.359
$\sin^2 \theta_{23}$	0.42	0.34-0.64
$\sin^2 \theta_{13}$	0.021	0.001-0.044

TABLE I: The best-fit and 3σ values for the mass squared differences and mixing angles from [28].

σ ranges of the mass squared differences and mixing angles from global analysis of oscillation data are depicted in Table I. Recently the evidence of non-zero θ_{13} at more than 5σ were reported by the Daya-Bay [26] and RENO [27] experiments with best-fit values and 3σ ranges of $\sin^2 \theta_{13}$ as,

$$\text{Daya} - \text{Bay} : \quad \sin^2 \theta_{13} = 0.023 \text{ (best - fit); } 0.009 - 0.037 \text{ (} 3\sigma \text{ range)}$$

$$\text{RENO} : \quad \sin^2 \theta_{13} = 0.026 \text{ (best - fit); } 0.015 - 0.041 \text{ (} 3\sigma \text{ range)}.$$

Thus there is an increase in the lower limit and a decrease in the upper limit of θ_{13} as compared to the values given in Table I.

With the parametrization of the mixing matrix in Eq.(12) the effective mass $|m_\nu^{ee}|$ becomes,

$$|m_\nu^{ee}| = |m_1 c_{12}^2 c_{13}^2 + m_2 s_{12}^2 c_{13}^2 e^{2i\alpha_2} + m_3 s_{13}^2 e^{2i\alpha_3}|. \quad (13)$$

The effective mass assumes different values depending on whether the neutrino mass states follow normal or inverted hierarchy or they are quasi-degenerate. Where,

- Normal hierarchy (NH) refers to the arrangement which corresponds to $m_1 < m_2 \ll m_3$ with

$$m_2 = \sqrt{m_1^2 + \Delta m_{\text{sol}}^2}, \quad m_3 = \sqrt{m_1^2 + \Delta m_{\text{atm}}^2 + \Delta m_{\text{sol}}^2}. \quad (14)$$

- Inverted hierarchy (IH) implies $m_3 \ll m_1 \sim m_2$ with

$$m_1 = \sqrt{m_3^2 + \Delta m_{\text{atm}}^2}, \quad m_2 = \sqrt{m_3^2 + \Delta m_{\text{sol}}^2 + \Delta m_{\text{atm}}^2}. \quad (15)$$

- Quasi degenerate neutrinos correspond to $m_1 \approx m_2 \approx m_3 \gg \sqrt{\Delta m_{\text{atm}}^2}$.

Fig.(1) displays the effective mass governing $0\nu\beta\beta$ as a function of the lowest mass scale in the standard three generation picture for various mass schemes. The gray (lighter) band for NH corresponds to varying the parameters in their 3σ range as given in Table I whereas the black band corresponds to the best-fit parameters. In both figures the Majorana phases are varied between 0 to 2π . The figure in the left panel is for 3σ ranges of $\sin^2 \theta_{13}$ from Table I whereas the plots in the right panel use 3σ range of $\sin^2 \theta_{13}$ as measured by the Daya-Bay experiment. Below we discuss briefly the limiting values of effective mass for different mass schemes. We also comment on the impact of the Daya-Bay and RENO measurement of θ_{13} on the effective mass. The different limits of Eq.(13) depend on the relative magnitudes of m_1 , $\sqrt{\Delta m_{\text{sol}}^2} \sim 9 \times 10^{-3}$ eV, $\sqrt{\Delta m_{\text{atm}}^2} \sim 0.05$ eV. Of special importance is the mass ratio $r = \left| \frac{\Delta m_{\text{sol}}^2}{\Delta m_{\text{atm}}^2} \right|$. In Table II we put the 3σ ranges of different parameters and their combinations that will be relevant for our discussion.

	\sqrt{r}	$\sqrt{r} \sin^2 \theta_{12}$	$\sqrt{r} \cos 2\theta_{12}$	$\tan^2 \theta_{13}$	$\sqrt{r} \tan^2 \theta_{13}$
Maximum	0.2	.072	.096	.046 (.037)	0.009 (0.007)
Minimum	0.16	.042	.046	.001 (.009)	0.0001 (0.002)

TABLE II: The 3σ ranges of different combinations of oscillation parameters relevant for understanding the behaviour of the effective mass in different limits. The values in the parentheses in the last two columns are from Daya-Bay results.

A. Normal Hierarchy

In the strictly hierarchical regime ($m_1 \ll m_2 \approx \sqrt{\Delta m_{\text{sol}}^2} \ll m_3 \approx \sqrt{\Delta m_{\text{atm}}^2}$) the effective mass due to the light neutrinos, can be approximated by

$$|m_\nu^{ee}|_{NH} = \sqrt{\Delta m_{\text{atm}}^2} \left| \sqrt{r} s_{12}^2 c_{13}^2 e^{2i\alpha_2} + s_{13}^2 e^{2i\alpha_3} \right|. \quad (16)$$

The maximum value of this corresponds to the phase-choice $\alpha_2 = \alpha_3 = 0$ while the minimum occurs for $\alpha_2 = 0, \alpha_3 = \pi$. Since the upper limit of $\sin^2 \theta_{13}$ from Daya-Bay measurement is less than the upper limit of $\sin^2 \theta_{13}$ given in Table I the maximum value of $|m_\nu^{ee}|_{NH}$ becomes lower and minimum value becomes higher in the second panel of Fig.(1). In this limit both the terms can be comparable resulting in complete cancellation if the following condition is

fulfilled:

$$\sqrt{r} \sin^2 \theta_{12} = \tan^2 \theta_{13}. \quad (17)$$

Comparing the columns 3 and 5 of Table II we see that using the range of θ_{13} measured by Daya-Bay experiment, complete cancellation is not satisfied exactly. This condition changes as we increase m_1 and approach the limit of partial hierarchy: $m_1 \approx m_2 \approx \sqrt{\Delta m_{sol}^2} \ll m_3 \approx \sqrt{\Delta m_{atm}^2}$ the minimum value for the above expression correspond to the phase choices $\alpha_2 = \alpha_3 = \pi/2$, The condition for complete cancellation now alters to,

$$\sqrt{r} \cos 2\theta_{12} = \tan^2 \theta_{13}. \quad (18)$$

From the 3σ ranges of parameters in Table II we see that in this region complete cancellation takes place in both panels in Fig.(1) resulting in very low values of the effective mass. Thus, maximum and minimum values of the effective mass is sensitive to the value of θ_{13} in this region.

B. Inverted Hierarchy

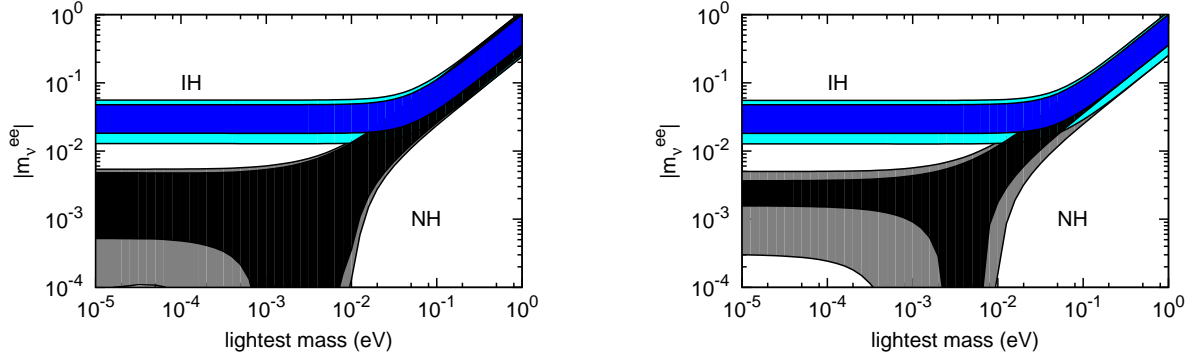


FIG. 1: The canonical contribution from light neutrino mass with θ_{13} from [26](left) and from [28](right) to the neutrinoless double beta decay.

In the inverted hierarchical case, $m_1 \approx m_2 \gg m_3$ and for the smaller values of m_3 such that $m_3 \ll \sqrt{\Delta m_{atm}^2}$, $m_2 \approx m_1 \approx \sqrt{\Delta m_{atm}^2}$ the effective mass is given as

$$|m_\nu^{ee}|_{IH} = \sqrt{\Delta m_{atm}^2} (c_{12}^2 c_{13}^2 + s_{12}^2 c_{13}^2 e^{2i\alpha_2}). \quad (19)$$

$|m_\nu^{ee}|$ in this region is independent of m_3 and is bounded from above and below by a maximum and minimum value given by

$$|m_\nu^{ee}|_{max} = c_{13}^2 \sqrt{\Delta m_{atm}^2} \quad (\alpha_2 = 0, \pi) \quad (20)$$

$$|m_\nu^{ee}|_{min} = c_{13}^2 \cos 2\theta_{12} \sqrt{\Delta m_{atm}^2} \quad (\alpha_2 = \pi/2, 3\pi/2). \quad (21)$$

As m_3 approaches $\sqrt{\Delta m_{atm}^2}$, the other two masses can be approximated as, $m_2 \approx m_1 \approx \sqrt{2\Delta m_{atm}^2}$ and the effective mass becomes,

$$|m_\nu^{ee}|_{IH} = \sqrt{\Delta m_{atm}^2} \left(\sqrt{2} c_{13}^2 (c_{12}^2 + s_{12}^2 e^{2i\alpha_2}) + s_{13}^2 e^{2i\alpha_3} \right) \quad (22)$$

The maximum value of this corresponds to $\alpha_2 = \alpha_3 = 0$ and complete cancellation cannot take place. For still higher values of m_3 one transcends into the quasi degenerate limit.

C. Quasi Degenerate

In this regime the effective mass (for both normal and inverted ordering) is given as

$$|m_\nu^{ee}|_{QD} = m_0 |c_{12}^2 c_{13}^2 + s_{12}^2 c_{13}^2 e^{2i\alpha_2} + s_{13}^2 e^{2i\alpha_3}|, \quad (23)$$

where $m_0 \approx m_1 \approx m_2 \approx m_3$. Thus the effective mass due to the light neutrinos increases linearly as the common mass scale.

These well known features are reflected in Fig.(1). In particular the Fig.(1)(right) shows the impact of Daya-Bay observations. This impact seems to be nominal for IH. For NH it affects the low m_1 region.

IV. NEUTRINOLESS DOUBLE- β DECAY IN MINIMAL LR MODEL

If we consider the left-right symmetric model then there can be several additional diagrams contributing to this process [29]:

- (a) The right-handed current mediated by W_R can contribute to the process through the exchange of the heavy neutrino N_R ,
- (b) There can also be light-heavy neutrino mixing diagrams with amplitudes $\sim m_D/M_R$,
- (c) Additional diagrams can also arise due to $W_L - W_R$ mixing,
- (d) Apart from above there can be additional contributions due to the charged Higgs fields

belonging to the bidoublet and triplet representations.

We discuss the relevant Lagrangian and the amplitudes of various contributions in the Appendix. In what follows we separate our discussion into three parts (i) type-I dominance, (ii) type-II dominance and (iii) contribution from the triplet Higgs.

A. Type-I Dominance

In this section we discuss the $0\nu\beta\beta$ for LR symmetric models with type-I dominance with the W_R mass to be of the order of $\sim \text{TeV}$. The type-I term dominates if one takes $v_L = 0$. It is well known that the minimization of the potential in the LR symmetric model admits this possibility [30]. Then,

$$m_\nu^{light} = \frac{v^2}{v_R} y_D^T f^{-1} y_D, \quad (24)$$

$$m_R^{heavy} = f v_R, \quad (25)$$

where we have used $f_L = f_R = f$. If U_R is the unitary matrix that diagonalizes the heavy neutrino mass matrix m_R , then since v_R is a constant, the same U_R will also diagonalize f and therefore one can write

$$\Rightarrow f^{-1} = U_R^T (f^{dia})^{-1} U_R \quad (26)$$

Using this in the expression for m_ν^{light} one gets

$$\begin{aligned} U_R m_\nu U_R^T &= \frac{v^2}{v_R} U_R y_D^T U_R^T (f^{dia})^{-1} U_R y_D U_R^T \\ &= \frac{v^2}{v_R} y_D^T (f^{dia})^{-1} y_D \\ &= m_\nu^{dia} \end{aligned} \quad (27)$$

Where in the last line we have used $U_R y_D U_R^T = y_D$ which is of course a special choice to establish a simplified relation between light and heavy neutrino masses.

The above choice implies

$$U_L = U_R, \quad \text{and} \quad m_i \propto \frac{1}{M_i}. \quad (28)$$

The proportionality constant depends on the vev's of Higgs which is generation independent and also y_D which depends on i .

In order to write the expression for the half-life of $0\nu\beta\beta$ process one needs to know the relative contribution of different terms. Since the mass of the right-handed gauge boson is

related to the scale v_R , therefore for $M_{W_R} \sim \text{TeV}$ the mass of the right-handed neutrinos is also $\sim \text{TeV}$ scale. However since $m_D^2/M_R \sim 0.01 - 0.1 \text{ eV}$, m_D^2 should be $\sim 10^{10} - 10^{11} \text{ eV}^2$ which in turn implies $m_D/M_R \sim 10^{-5} - 10^{-6}$. Thus light-heavy mixing remains very small unless one does some fine tuning of the Yukawa textures [31]. In this approximation, the only processes that contribute to the $0\nu\beta\beta$ are the W_L mediated diagrams through light neutrino exchange and W_R mediated diagrams through heavy neutrino exchange, as discussed in the Appendix.

The Δ_R mediated diagram can in principle contribute for W_R mass around TeV scale. But invoking the constraint from Lepton Flavour Violating(LFV) decays it is seen that for majority of the parameter space $M_N/M_\Delta < 0.1$ and hence the Δ_R contribution can be suppressed [16, 32]. We will comment later on this scenario for the case $M_N \approx M_\Delta$.

Under the above approximations the time-period for $0\nu\beta\beta$ process can be written as:

$$\frac{\Gamma_{0\nu\beta\beta}}{\ln 2} = G \left(|\mathcal{M}_\nu|^2 |\eta_L|^2 + |\mathcal{M}_N|^2 |\eta_R|^2 \right), \quad (29)$$

where

$$|\eta_L| = \frac{|U_{Lei}^2 m_i|}{m_e} = m_\nu^{ee}/m_e, \quad (30)$$

$$|\eta_R| = \frac{M_{W_L}^4}{M_{W_R}^4} |(U_R)_{ei}^2 m_p/M_i|. \quad (31)$$

The interference diagram between these two terms is helicity suppressed being proportional to the electron mass. The contribution from the neutrino propagator term to the amplitude is $\sim \frac{m_i}{p^2 - m_i^2}$. The different dependence on the masses for the left and right sector come since the exchanged momentum p satisfies $m_i \ll p \ll M_i$. \mathcal{M}_ν and \mathcal{M}_N are the nuclear matrix elements corresponding to light and heavy neutrino exchange respectively.

Eq.(29) can be expressed as,

$$\frac{\Gamma_{0\nu\beta\beta}}{\ln 2} = \frac{G |\mathcal{M}_\nu|^2}{m_e^2} |m_{ee}^{\text{eff}}|^2. \quad (32)$$

Therefore the effective neutrino mass governing neutrinoless double beta decay is

$$|m_{ee}^{\text{eff}}|^2 = |m_\nu^{ee}|^2 + |M_N^{ee}|^2, \quad (33)$$

where,

$$|M_N^{ee}| = \left| \langle p \rangle^2 \frac{M_{W_L}^4}{M_{W_R}^4} \frac{(U_R)_{ei}^2}{M_i} \right|, \quad (34)$$

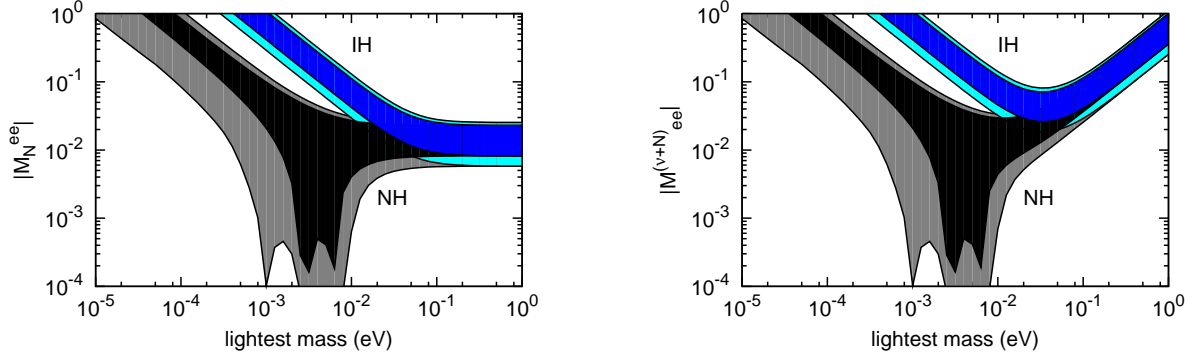


FIG. 2: Contributions to effective mass $|M_N^{ee}|$ from only right-handed neutrino (left) and the total contributions $|M_{N+\nu}^{ee}|$ from left+right handed neutrinos (right) in case of type-I dominance.

is the contribution to effective mass from the right-handed current. In the above expression, $|\langle p^2 \rangle| = |m_e m_p \mathcal{M}_N / \mathcal{M}_\nu|$.

In Fig.(2) we show the effective mass in LR symmetric models under type-I seesaw dominance. The right-handed neutrino mass M_{W_R} is taken as 3.5 TeV. The mass of the heaviest right-handed neutrino is 500 GeV. The allowed value of p is in the range $\sim (100-200)$ MeV. In our analysis we have adopted $p = 180$ MeV, $M_{W_L} = 81$ GeV, $C_N = p^2 (\frac{M_W}{M_{W_R}})^4 \sim 10^{10} \text{ eV}^2$. The left panel shows the contribution from the right-handed current whereas the right panel shows the total contribution. For these plots we have used the 3σ range of $\sin^2 \theta_{13}$ as from the Daya-Bay results. Other oscillation parameters are varied in their 3σ ranges as given in Table I. In order to understand the various features of these diagrams and the interplay of the different contributions we examine the expression of effective mass in various limits for NH, IH and QD neutrinos.

1. Normal Hierarchy

In normal hierarchy regime as m_1 is the lightest neutrino, the heaviest RH neutrino will be M_1 as seen from the above Eq.(28). The RH neutrino masses M_2, M_3 can therefore be expressed in terms of the heaviest RH mass as

$$\frac{M_2}{M_1} = \frac{m_1}{m_2}, \quad (35)$$

$$\frac{M_3}{M_1} = \frac{m_1}{m_3}. \quad (36)$$

To find the above simplified relation between light and heavy neutrino masses we consider the eigenvalues of y_D matrix to be degenerate, i.e., $y_{D1} = y_{D2} = y_{D3}$. We get the effective mass due to contribution from the RH neutrinos using above equation Eq.(28)

$$\left| M_N^{ee} \right| = C_N \left| \frac{c_{12}^2 c_{13}^2}{M_1} + \frac{s_{12}^2 c_{13}^2}{M_2} e^{2i\alpha_2} + \frac{s_{13}^2}{M_3} e^{2i\alpha_3} \right| \quad (37)$$

$$= \frac{C_N}{M_1} \left| c_{12}^2 c_{13}^2 + \frac{m_2}{m_1} s_{12}^2 c_{13}^2 e^{2i\alpha_2} + \frac{m_3}{m_1} s_{13}^2 e^{2i\alpha_3} \right|. \quad (38)$$

In the limit when $m_1 \ll m_2 \simeq \sqrt{\Delta m_{sol}^2} \ll m_3 \simeq \sqrt{\Delta m_{atm}^2}$ the above term reads as

$$\left| M_N^{ee} \right| = \frac{C_N}{M_1} \left| c_{12}^2 c_{13}^2 + \frac{\sqrt{\Delta m_{sol}^2}}{m_1} s_{12}^2 c_{13}^2 e^{2i\alpha_2} + \frac{\sqrt{\Delta m_{atm}^2}}{m_1} s_{13}^2 e^{2i\alpha_3} \right|. \quad (39)$$

This gives a steep increase in effective mass as we go towards the smaller m_1 . This is in complete contrast with the effective mass term $|m_\nu^{ee}|$. In this region therefore the total contribution is dominated by the RH sector.

As m_1 increases and to $m_1 \simeq m_2 \simeq \sqrt{\Delta m_{sol}^2} \ll m_3 \simeq \sqrt{\Delta m_{atm}^2}$. The effective mass term for the right-handed neutrino is

$$\left| M_N^{ee} \right| = \frac{C_N}{M_1} \left| c_{13}^2 (c_{12}^2 + s_{12}^2 e^{2i\alpha_2}) + \sqrt{\frac{1}{r}} s_{13}^2 e^{2i\alpha_3} \right|, \quad (40)$$

for $\alpha_2 = 0, \alpha_3 = \pi/2$ we get

$$\sqrt{r} = \tan^2 \theta_{13}, \quad (41)$$

for $\alpha_2 = \pi/2, \alpha_3 = \pi/2$

$$\sqrt{r} \cos 2\theta_{12} = \tan^2 \theta_{13}. \quad (42)$$

The first condition cannot be satisfied by the present values of the oscillation parameters as is evident from Table II. But the second condition is same as what we have got for the light neutrino case and cancellations can occur in this range for the right handed sector as well.

2. *Inverted Hierarchy*

In inverted hierarchy regime the lightest neutrino is the m_3 . Therefore the heaviest RH neutrino in this ordering will be M_3 as is evident from Eq.(28). The RH neutrino masses M_1, M_2 can therefore be expressed in terms of the heaviest RH Mass as

$$\frac{M_2}{M_3} = \frac{m_3}{m_2}, \quad (43)$$

$$\frac{M_1}{M_3} = \frac{m_3}{m_1}. \quad (44)$$

The heavy neutrino contribution to the effective mass is given by

$$\left| M_N^{ee} \right| = \frac{C_N}{M_3} \left| \frac{m_1}{m_3} c_{12}^2 c_{13}^2 + \frac{m_2}{m_3} s_{12}^2 c_{13}^2 e^{2i\alpha_2} + s_{13}^2 e^{2i\alpha_3} \right|. \quad (45)$$

In the limit $m_3 \ll \sqrt{\Delta m_{atm}^2} \ll m_2 \simeq m_1 \simeq \sqrt{\Delta m_{atm}^2}$ the RH contribution to the effective mass is

$$\left| M_N^{ee} \right| = \frac{C_N}{M_3} \left| \frac{\sqrt{\Delta m_{atm}^2}}{m_3} c_{12}^2 c_{13}^2 + \frac{\sqrt{\Delta m_{atm}^2}}{m_3} s_{12}^2 c_{13}^2 e^{2i\alpha_2} + s_{13}^2 e^{2i\alpha_3} \right|. \quad (46)$$

In the above equation the first and the second terms dominate and the effective mass reveal an increase with decrease in m_3 for the IH case as well.

3. *Quasi Degenerate*

In this limit $m_1 \simeq m_2 \simeq m_3$ which in turn implies $M_1 \approx M_2 \approx M_3 \approx M_0$. The RH contribution to the effective mass can be expressed as,

$$\left| M_N^{ee} \right| = \frac{C_N}{M_0} \left| c_{12}^2 c_{13}^2 + s_{12}^2 c_{13}^2 e^{2i\alpha_2} + s_{13}^2 e^{2i\alpha_3} \right|, \quad (47)$$

Thus the effective mass is independent of the light neutrino mass and remains constant. Since the light neutrino contribution in this limit increases with the mass scale this part dominates in the total contribution resulting in an overall increase in the total effective mass with increasing m_1 .

B. *Type-II dominance*

Type-II dominance implies the Dirac term connecting the left- and the right- handed states is negligibly small as compared to the type-I term. In this limit we can write

$$\begin{aligned} m_\nu^{light} &= f_L v_L, \\ m_R^{heavy} &= f_R v_R. \end{aligned}$$

Denoting the matrices diagonalizing m_ν^{light} and m_R^{heavy} as U_L and U_R respectively, one can have two possibilities

$$f_L = f_R \implies U_L = U_R, \quad (48)$$

$$f_L = f_R^* \implies U_L = U_R^*. \quad (49)$$

Using Eq.(48) relation between the light and heavy masses as,

$$m_i \propto M_i. \quad (50)$$

Note that the proportionality constant in this case is v_L/v_R and is independent of the generation index i . In this case, using Eq.(50) one can relate the heavy neutrino mass ratios to those of light neutrinos as,

$$\frac{M_1}{M_3} = \frac{m_1}{m_3}$$

and $\frac{M_2}{M_3} = \frac{m_2}{m_3}.$

The dominant contributions to $0\nu\beta\beta$ come from two diagrams one via exchange of the light neutrinos and another via exchange of the heavy neutrinos. The charged Higgs diagram as well as the $W_L - W_R$ mixing diagrams can be neglected in this limit [16, 32]. Therefore the time-period and the effective mass is given by the same expression as in Eq. (33) for the type-I case.

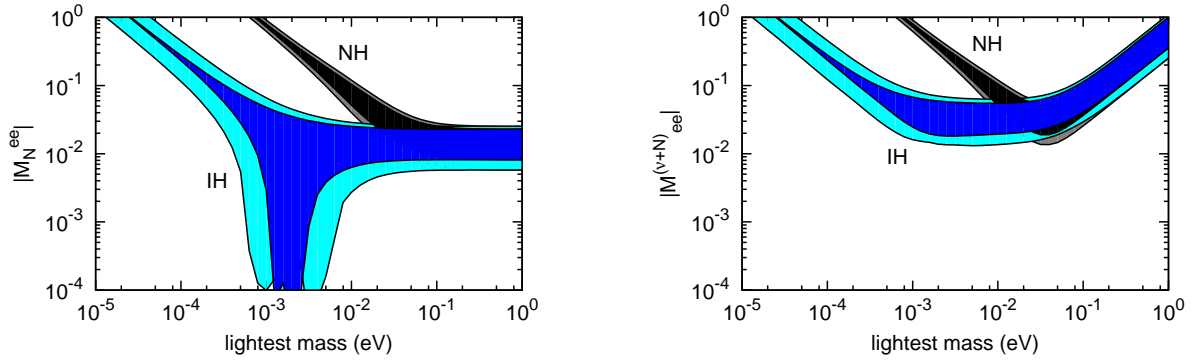


FIG. 3: Contributions to effective mass $|M_N^{ee}|$ from only right-handed neutrino (left) and the total contributions $|M_{N+\nu}^{ee}|$ from left+right-handed neutrinos (right) in case of type-II dominance.

Below we discuss the contributions to the effective mass from the heavy sector and the interplay of the light and heavy contribution to the total effective mass for NH, IH and QD neutrinos.

1. Normal Hierarchy

In order to examine the behavior of the RH current, we fix M_3 (say around 500 GeV). Once M_3 is fixed, the other heavy neutrino masses can be expressed in terms of light neutrino data and M_3 . Also, one can assume that $U_R = U_L$. We have assigned the same values of p , M_{W_L} and M_{W_R} as mentioned in type-I dominance case. The expression for M_N^{ee} now becomes,

$$|M_N^{ee}|_{NH} = \frac{C_N}{M_3} \left| c_{12}^2 c_{13}^2 \frac{m_3}{m_1} + s_{12}^2 c_{13}^2 e^{2i\alpha_2} \frac{m_3}{m_2} + s_{13}^2 e^{2i\alpha_3} \right|. \quad (51)$$

In the limit of strict hierarchy: $m_1 \ll m_2 \approx \sqrt{\Delta m_{sol}^2} \ll m_3 \approx \sqrt{\Delta m_{atm}^2}$, the heavy neutrino contribution M_N^{ee} can be written as,

$$|M_N^{ee}| = \frac{C_N}{M_3} \left| c_{12}^2 c_{13}^2 \frac{\sqrt{\Delta m_{atm}^2}}{m_1} + s_{12}^2 c_{13}^2 e^{2i\alpha_2} \frac{1}{\sqrt{r}} + s_{13}^2 e^{2i\alpha_3} \right|. \quad (52)$$

For smaller values of m_1 , the first term dominates showing a steep rise in the effective mass parameter as we go towards smaller m_1 . This is shown in Fig.(3) in the left panel. This behaviour is in complete contrast with the effective mass $|m_\nu^{ee}|$ due to the left-handed current which can become vanishingly small due to complete cancellation between the various contributions. Since the heavy neutrino contribution is much higher, in the total effective mass, this plays the dominant role and hence the total effective mass shows a sharp decrease with m_1 in this regime as is seen from the right panel of Fig.(3). In the limit $m_1 \approx m_2 \approx \sqrt{\Delta m_{sol}^2} \ll m_3 \approx \sqrt{\Delta m_{atm}^2}$

$$|M_N^{ee}|_{NH} = \frac{C_N}{M_3} \left| \frac{c_{13}^2}{\sqrt{r}} \left(c_{12}^2 + s_{12}^2 e^{2i\alpha_2} \right) + s_{13}^2 e^{2i\alpha_3} \right|. \quad (53)$$

The minimum value correspond to $\alpha_2 = 0, \alpha_3 = \pi/2$ or $\alpha_2 = \alpha_3 = \pi/2$. Complete cancellation would require

$$\sqrt{r} \tan^2 \theta_{13} = c_{12}^2 \pm s_{12}^2. \quad (54)$$

The left-hand side is a product of two small numbers and for the current 3σ range of parameters the above condition is not satisfied (cf. Table II) and hence complete cancellation cannot occur in this limit as well. As a result in the total effective mass the heavy neutrino contribution dominates in this regime.

2. *Inverted Hierarchy*

In order to examine the analytical behavior for RH current for IH we fix the highest mass state M_2 around 500 GeV. Once M_2 is fixed, the other heavy neutrino masses can be expressed in terms of light neutrino masses and M_2 as

$$\frac{M_1}{M_2} = \frac{m_1}{m_2},$$

and $\frac{M_3}{M_2} = \frac{m_3}{m_2}.$

Now the expression for M_N^{ee} becomes

$$|M_N^{ee}|_{IH} = \frac{C_N}{M_2} \left| c_{12}^2 c_{13}^2 \frac{m_2}{m_1} + s_{12}^2 c_{13}^2 e^{2i\alpha_2} + s_{13}^2 e^{2i\alpha_3} \frac{m_2}{m_3} \right|. \quad (55)$$

For the smaller values of $m_3 \ll \sqrt{\Delta m_{\text{atm}}^2} \approx m_2 \approx m_1$,

$$|M_N^{ee}|_{IH} = \frac{C_N}{M_2} \left| c_{12}^2 c_{13}^2 + s_{12}^2 c_{13}^2 e^{2i\alpha_2} + s_{13}^2 \frac{\sqrt{\Delta m_{\text{atm}}^2}}{m_3} e^{2i\alpha_3} \right|. \quad (56)$$

It is clear from the above expression that for smaller values of m_3 the absolute value of effective neutrino mass varies with the lightest neutrino mass as $1/m_3$. Since the contribution m_ν^{ee} in this region is smaller the total contribution is dominated by $|M_N^{ee}|$ and show the $1/m_3$ decrease. As m_3 increases the third term starts becoming smaller and there can be complete cancellations. There are two end points corresponding to $\alpha_2 = 0, \pi$, $\alpha_3 = \pi/2$ and $\alpha_2 = \pi/2$, $\alpha_3 = \pi/2$) where M_N^{ee} acquires the minimum value. This leads to the following two conditions for cancellation region

$$m_3 = \sqrt{\Delta m_{\text{atm}}^2} \tan^2 \theta_{13} \quad (57)$$

$$m_3 = \sqrt{\Delta m_{\text{atm}}^2} \tan^2 \theta_{13} \cos 2\theta_{12}. \quad (58)$$

The Eq.(57) gives $m_3 \sim 6 \times 10^{-5} - 2 \times 10^{-3}$ eV for the 3σ ranges of parameters from Table I (but $\sin^2 \theta_{13}$ from Daya-Bay) while Eq.(58) gives $m_3 = 1.7 \times 10^{-4} - 9.2 \times 10^{-4}$. These ranges define the values of m_3 for which complete cancellation can occur. This is reflected in Fig(3). However when we consider the total effective mass then the contribution from $|m_\nu^{ee}|$ dominates and enforces a lower limit on this. Thus vanishing effective mass is not a possibility for both hierarchies in presence of right-handed currents and type-II seesaw dominance.

3. *Quasi Degenerate*

Quasi-degeneracy in the light neutrino masses also implies $M_1 \approx M_2 \approx M_3 \approx M_0$ for the heavy neutrinos. In this regime the heavy neutrino contribution to the effective mass is

$$|M_N^{ee}|_{QD} = \frac{C_N}{M_0} \left| c_{12}^2 c_{13}^2 + s_{12}^2 c_{13}^2 e^{2i\alpha_2} + s_{13}^2 e^{2i\alpha_3} \right|. \quad (59)$$

This is independent of the lightest neutrino mass and hence value of $|M_N^{ee}|$ remains constant with increasing value of the common mass scale. The overall behaviour of the total effective mass in this regime is therefore controlled by the lighter neutrino contribution and increases with increasing mass.

C. The contribution from the triplet Higgs

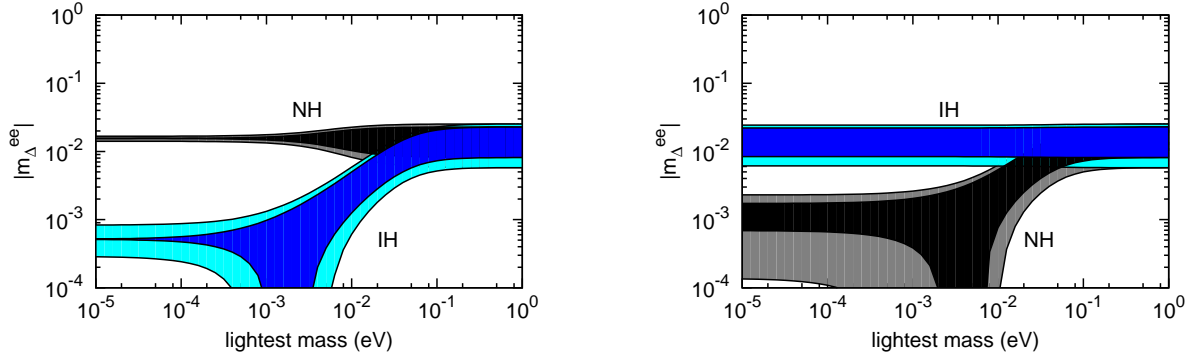


FIG. 4: The contribution to the effective mass from right-handed triplet diagrams for type-I seesaw (left) and type-II seesaw (right) diagrams respectively.

The Majorana masses of light and heavy neutrinos comes naturally in left-right model because of the two triplets $\Delta_{L,R}$. As discussed in the appendix the contribution from Δ_L is much suppressed as compared to the dominant contributions. However the magnitude of the Δ_R contribution is controlled by the factor M_i/M_{Δ_R} . In the earlier sections we have not included the contribution due to the triplet Higgs diagrams under the assumption $M_i/M_{\Delta_R} < 0.1$, which is obtained from LFV processes [16]. However, this approximation though valid in a large part of the parameter space there are some allowed mixing parameters for which this ratio can be higher [16]. In that case the contribution from this diagram

needs to be included. In order to assess the impact of this contribution below we discuss the effective mass due to the Δ_R^{--} diagram (see Appendix) in the limit $M_{heaviest} = M_{\Delta_R}$.

Comparing with the light neutrino exchange diagram, the Δ_R exchange diagram (which is shown in Fig.(4)) gives the effective mass as,

$$|m_{\Delta}^{ee}| = \left| p^2 \frac{M_{W_L}^4}{M_{W_R}^4} \frac{2 M_N}{M_{\Delta_R}^2} \right|. \quad (60)$$

In Fig.(4) we plot the contribution of the effective mass due to the triplets, $|m_{\Delta}^{ee}|$, for type-I and type-II seesaw model. The figure shows that for the type-I case and NH the triplet contribution can dominate over the light and heavy neutrino contributions in the cancellation regime. However for the other cases, i.e., type-I and IH and type-II (both NH and IH), the triplet contributions are always lower than the other two contributions.

Note that in generating the above plots we have varied the mixing parameters in their 3σ range and all phases between 0 to 2π . However, note that the approximation $M_{heaviest} = M_{\Delta_R}$ may not be valid for all the parameter values and a more accurate analysis would require a correlated study of LFV and $0\nu\beta\beta$ which is beyond the scope of this paper.

V. CONCLUSIONS

Neutrino oscillation experiments have already provided us the first signature of physics beyond the Standard Model. At the present juncture the quest for new physics has also got an unprecedented momentum because of LHC. A natural question is whether the origin of neutrino mass can be probed at LHC and it is hoped that the mist around the TeV scale may be uplifted by complimentary information from neutrino and collider experiments. Among neutrino experiments, observation of $0\nu\beta\beta$ would signify lepton number violation and Majorana nature of neutrino mass. However as is well known $0\nu\beta\beta$ can also occur in many other scenarios and hence the specific nature of new physics may remain to be ascertained and in such situations LHC and LFV processes may provide complimentary information. This interrelation of $0\nu\beta\beta$ with LHC and LFV processes makes it a very engrossing and interesting topic of research at the present juncture. In this paper we scrutinize the implications of $0\nu\beta\beta$ in left-right symmetric models with right handed gauge bosons and neutrinos around TeV scale. We analyze the effective mass within the frameworks of type-I and type-II dominance cases. For the latter case the hierarchy in heavy sector is same as that in

the light sector. For the former, we use some simplifying assumptions for relating the mass hierarchy of the heavy sector with that of the light sector. We identify the situations where the dominant contribution can come from the right-handed sectors. The various parameters chosen for our numerical work are consistent with constraints from neutrino oscillation experiments as well as results from LHC and Lepton Flavour Violation. In the light of recent announcement on the measurement of the third leptonic mixing angle from reactor experiments Daya Bay and RENO we have discussed the impacts of the measured value of θ_{13} on the behaviour of effective mass. The regime of cancellations between different contributions to $0\nu\beta\beta$ and their compatibility with present oscillation data are also presented. We note that in both type-I, and type-II dominance situations the right-handed contribution can override the left-handed ones for smaller values of the lightest neutrino mass while in the quasi-degenerate regime the light neutrino contribution dominates. For the type-II dominance case the predictions for NH and IH including right-handed currents are in stark contrast compared to the case of only left-handed currents. Thus, signature of TeV scale LR symmetry at LHC may present a completely altered interpretation for effective mass governing $0\nu\beta\beta$ process.

VI. APPENDIX

The canonical type-I seesaw in left-right symmetric model is automated through the presence of the bidoublet that generates the Dirac coupling, m_D between left- and right-handed neutrinos while the type-II seesaw term m_L is due to the presence of the triplet Higgs. The relevant Lagrangian reads as

$$\mathcal{L} = -\frac{1}{2} \left(\overline{\nu'_L} \quad \overline{N'_R} \right) \mathcal{M} \begin{pmatrix} \nu'_L \\ N'_R \end{pmatrix} + \text{h.c.}, \quad (61)$$

where \mathcal{M} is the neutrino mass matrix which can be expressed as,

$$\mathcal{M} = \begin{pmatrix} m_L & m_D^T \\ m_D & M_R \end{pmatrix}_{6 \times 6}. \quad (62)$$

M_R is the Majorana mass matrix for the right-handed neutrinos which arises through the right handed triplet Higgs. The 6×6 neutrino mass matrix Eq.(62) can be diagonalized by

the 6×6 unitary matrix \mathcal{U} , defined as

$$\mathcal{U}^T \begin{pmatrix} 0 & m_D^T \\ m_D & M_R \end{pmatrix} \mathcal{U} = \begin{pmatrix} m^{diag} & 0 \\ 0 & M^{diag} \end{pmatrix} \quad \text{and} \quad \begin{pmatrix} \nu'_L \\ N_R^{c'} \end{pmatrix} = \mathcal{U} \begin{pmatrix} \nu_L \\ N_R^c \end{pmatrix}, \quad (63)$$

where $m^{diag} = \text{diag}(m_1, m_2, m_3)$ and $M^{diag} = \text{diag}(M_1, M_2, M_3)$ are the diagonal matrices with mass eigenvalues m_i and M_i ($i = 1, 2, 3$) for light and heavy neutrinos respectively. The primed and unprimed notations in the neutrino fields correspond to the flavour and mass eigenstates respectively. The complete unitary mixing matrix \mathcal{U} can be parametrized as a product of two matrices $\mathcal{U} = W U_\nu$ and can be expressed as [33]

$$\mathcal{U} = W U_\nu = \begin{pmatrix} (1 - \frac{1}{2} R R^\dagger) U'_L & R U'_R \\ -R^\dagger U'_L & (1 - \frac{1}{2} R^\dagger R) U'_R \end{pmatrix} = \begin{pmatrix} U_L & T \\ S & U_R \end{pmatrix}, \quad (64)$$

where W is the matrix which brings the 6×6 neutrino matrix, Eq.(62), in the block diagonal form and $U_\nu = \text{diag}(U'_L, U'_R)$ diagonalizes the left and right handed parts. The matrix R appearing in the Eq.(64) is $R = m_D^\dagger M_R^{-1*}$.

In left-right symmetric theories, the charged current interactions of leptons, in the flavour basis is given by

$$\mathcal{L}_{CC} = \frac{g}{\sqrt{2}} \sum_{\alpha=e,\mu,\tau} \left[\bar{\ell}_{\alpha L} \gamma_\mu \nu'_{\alpha L} W_L^\mu + \bar{\ell}_{\alpha R} \gamma_\mu N'_{\alpha R} W_R^\mu \right] + \text{h.c.} \quad (65)$$

Assuming a basis such that the charged-lepton mass matrix is diagonal, the above can be rewritten in the mass basis as,

$$\begin{aligned} \mathcal{L}_{CC} = & \frac{g}{\sqrt{2}} \sum_{\alpha=e,\mu,\tau} \sum_{i=1}^3 \left[\bar{\ell}_{\alpha L} \gamma_\mu \{ (U_L)_{\alpha i} \nu_{Li} + (T)_{\alpha i} N_{Ri}^c \} W_L^\mu \right. \\ & \left. + \bar{\ell}_{\alpha R} \gamma_\mu \{ (S)_{\alpha i}^* \nu_{Li}^c + (U_R)_{\alpha i}^* N_{Ri} \} W_R^\mu \right] + \text{h.c.} \end{aligned} \quad (66)$$

With this particular form of the charged current interaction, there are several additional diagrams to neutrinoless double beta decay. We have categorised them as: first that involves two W_L (the amplitude is denoted as \mathcal{A}^{LL}), the second that involves the exchange of two W_R gauge bosons \mathcal{A}^{RR} , and the third one that involves both W_L and W_R exchange at the same diagram \mathcal{A}^{LR} . The amplitudes of the contributions from the different diagrams are as indicated below :

$$(i) \quad \mathcal{A}_\nu^{LL} \propto \frac{1}{M_{W_L}^4} \frac{U_{Le_i}^2 m_i}{p^2} \quad (\text{Fig.6a}),$$

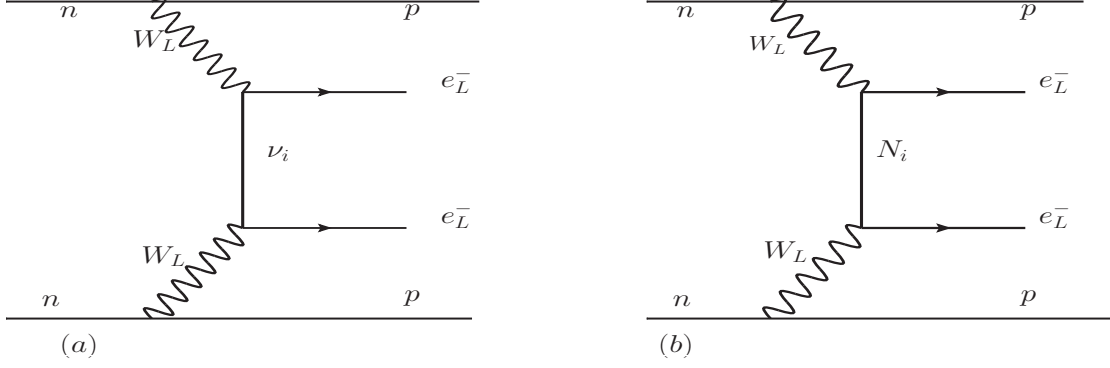


FIG. 5: Neutrinoless double beta decay contribution from light and heavy Majorana neutrino intermediate states from two W_L exchange.

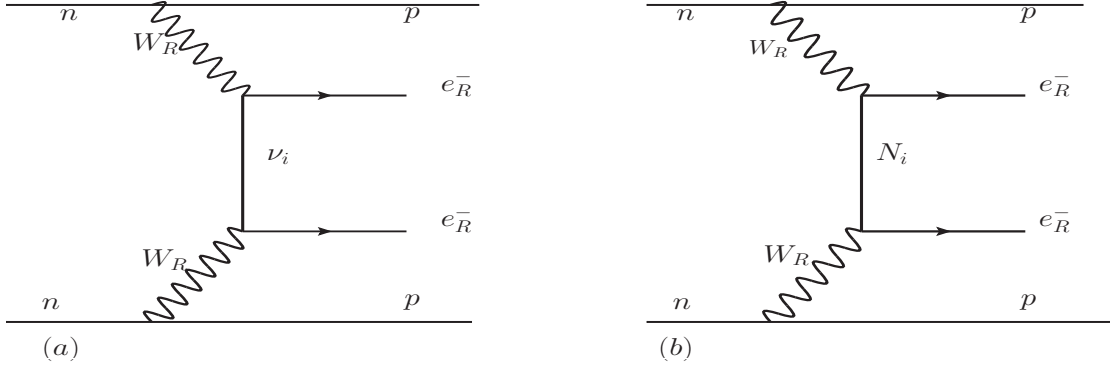


FIG. 6: Neutrinoless double beta decay contribution from light and heavy Majorana neutrinos from two W_R exchange.

$$(ii) \quad \mathcal{A}_N^{LL} \propto \frac{1}{M_{W_L}^4} \frac{T_{ei}^2}{M_i} \quad (\text{Fig. 6b}),$$

$$(iii) \quad \mathcal{A}_\nu^{RR} \propto \frac{1}{M_{W_R}^4} \frac{(S_{ei}^*)^2 m_i}{p^2} \quad (\text{Fig. 7a}),$$

$$(iv) \quad \mathcal{A}_N^{RR} \propto \frac{1}{M_{W_R}^4} \frac{(U_{Rei}^*)^2}{M_i} \quad (\text{Fig. 7b}),$$

$$(v) \quad \mathcal{A}_\nu^{LR} \propto \frac{1}{M_{W_L}^2 M_{W_R}^2} \frac{U_{Lei} S_{ei}^* m_i}{p^2} \quad (\text{Fig. 8a}),$$

$$(vi) \quad \mathcal{A}_N^{LR} \propto \frac{1}{M_{W_L}^2 M_{W_R}^2} \frac{T_{ei} U_{Rei}^*}{M_i} \quad (\text{Fig. 8b}).$$

There will also be the contribution from the triplet Higgs diagrams as given in Fig.(8).

The amplitudes for the diagrams can be written as,

$$(i) \quad \mathcal{A}_{\Delta_L}^{LL} \propto \frac{1}{M_{W_L}^4} \frac{1}{M_{\Delta_L}^2} f_L v_L,$$

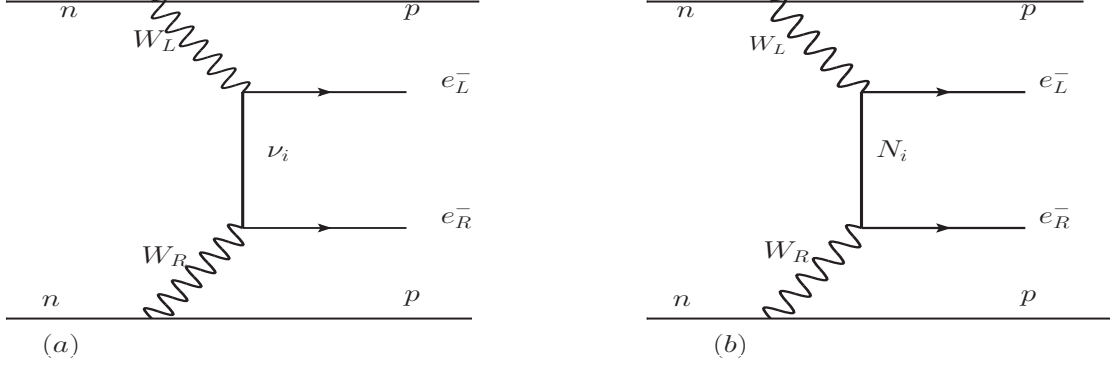


FIG. 7: Neutrinoless double beta decay contribution from light and heavy Majorana neutrino intermediate states from both left- and right- handed gauge bosons exchange at each vertices's.

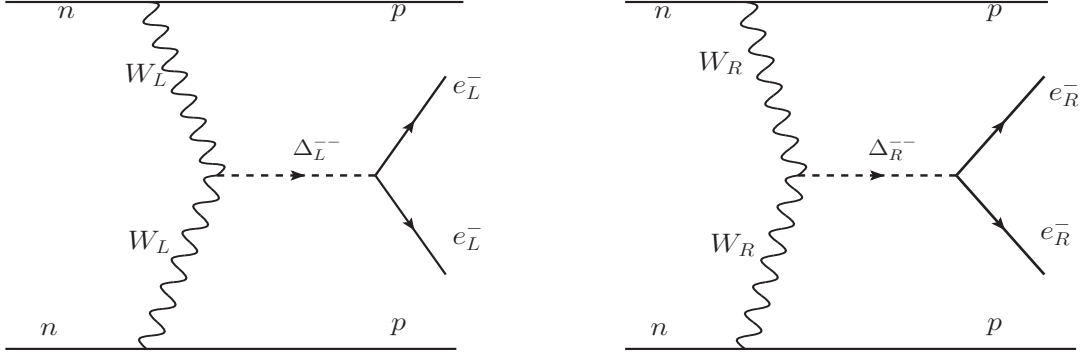


FIG. 8: Neutrinoless double beta decay contribution from the charged Higgs intermediate states from W_L and W_R exchange.

$$(ii) \quad \mathcal{A}_{\Delta_R}^{RR} \propto \frac{1}{M_{W_R}^4} \frac{1}{M_{\Delta_R}^2} f_R v_R.$$

The time period for $0\nu\beta\beta$ can be expressed as,

$$\frac{\Gamma_{0\nu\beta\beta}}{\ln 2} = G \left| \mathcal{M}_{\nu_L} \left(\eta_{\nu}^{LL} + \eta_{\nu}^{RR} + \eta_{\nu}^{LR} + \eta_{\Delta_L}^{LL} \right) + \mathcal{M}_{N_R} \left(\eta_N^{LL} + \eta_N^{RR} + \eta_N^{LR} + \eta_{\Delta_R}^{RR} \right) \right|^2, \quad (67)$$

where the dimensionless parameters η are defined as follows:

$$\eta_{\nu}^{LL} = \frac{U_{Lei}^2 m_i}{m_e}; \quad \eta_N^{LL} = \frac{T_{ei}^2 m_p}{M_i};$$

$$\eta_N^{RR} = \frac{M_{W_L}^4}{M_{W_R}^4} \frac{U_{Rei}^{*2} m_p}{M_i}; \quad \eta_{\nu}^{RR} = \frac{M_{W_L}^4}{M_{W_R}^4} \frac{S_{ei}^{*2} m_i}{m_e};$$

$$\eta_\nu^{LR} = \frac{M_{W_L}^2}{M_{W_R}^2} \frac{U_{Le i} S_{ei}^* m_i}{m_e}; \quad \eta_N^{LR} = \frac{M_{W_L}^2}{M_{W_R}^2} \frac{T_{ei} U_{Re i}^* m_p}{M_i};$$

$$\eta_{\Delta_L}^{LL} = \frac{U_{Le i}^2 m_i m_e}{M_{\Delta_L}^2}; \quad \eta_{\Delta_R}^{RR} = \frac{M_{W_L}^4}{M_{W_R}^4} \frac{U_{Re i}^2 M_i m_p}{M_{\Delta_R}^2}.$$

Inserting these the half-life is,

$$\begin{aligned} \frac{\Gamma_{0\nu\beta\beta}}{\ln 2} = G \frac{|\mathcal{M}_\nu|^2}{m_e^2} & \left| \left(U_{Le i}^2 m_i + p^2 \frac{T_{ei}^2}{M_i} + p^2 \frac{M_{W_L}^4}{M_{W_R}^4} \frac{U_{Re i}^{*2}}{M_i} \right. \right. \\ & + \frac{M_{W_L}^4}{M_{W_R}^4} S_{ei}^{*2} m_i + \frac{M_{W_L}^2}{M_{W_R}^2} U_{Le i} S_{ei}^* m_i + p^2 \frac{M_{W_L}^2}{M_{W_R}^2} \frac{T_{ei} U_{Re i}^*}{M_i} \\ & \left. \left. + \frac{U_{Le i}^2 m_i m_e^2}{M_{\Delta_L}^2} + p^2 \frac{M_{W_L}^4}{M_{W_R}^4} \frac{U_{Re i}^2 M_i}{M_{\Delta_R}^2} \right) \right|^2. \end{aligned} \quad (68)$$

In order to estimate the relative contributions of different terms note that our analysis is done $M_{W_R} \sim \text{TeV}$ and the heaviest right-handed neutrino is in the range $\sim 500 \text{ GeV}$. Since $m_\nu \sim m_D^2/M_R \sim 0.01 - 0.1 \text{ eV}$, $m_D \approx 10^5 \text{ eV}$ which in turn implies $m_D/M_R \sim 10^{-6} - 10^{-7} \sim T_{ei} \sim S_{ei}$ and $p^2 \sim 100 \text{ MeV}$. We also assume as illustrative values $M_{\Delta_L} \approx M_{\Delta_R} = 1 \text{ TeV}$.

Then the order of magnitude of the various terms in the above expression are as follows,

- $U_{ei}^2 m_i \sim m_i \sim 0.01 - 0.1$
- $p^2 \frac{T_{ei}^2}{M_i} \sim 10^{-8}$
- $p^2 \frac{M_{W_L}^4}{M_{W_R}^4} \frac{U_{Le i}^2}{M_i} \sim 0.01$
- $\frac{M_{W_L}^4}{M_{W_R}^4} S_{ei}^2 m_i \sim 10^{-18} m_i$
- $\frac{M_{W_L}^2}{M_{W_R}^2} S_{ei} U_{Re i} m_i \sim 10^{-9} m_i$
- $p^2 \frac{M_{W_L}^2}{M_{W_R}^2} \frac{T_{ei} U_{Re i}}{M_i} \sim 10^{-5}$
- $\frac{U_{Le i}^2 m_i m_e^2}{M_{\Delta_L}^2} \sim 10^{-13} m_i$
- $p^2 \frac{M_{W_L}^4}{M_{W_R}^4} \frac{U_{Re i}^2 M_i}{M_{\Delta_R}^2} \sim 10^{-5}.$

Therefore, in the approximation the non-unitarity of the mixing is small ($\sim 10^{-6}$) the dominant contribution will come from the left-handed current with light neutrino exchange and the right-handed current with the heavy neutrino exchange. However in cases where these contributions vanish due to some cancellations the effective mass can be generated from the diagram 8(b) since the same cancellation condition may not be operative in this term [34]. However because of T_{ei} term the overall contribution is still expected to be small for our choice of parameters. In our analysis we will therefore neglect this contribution. Also, if one assumes $M_i \approx M_{\Delta_R}$ then the Δ_R exchange diagram can also become significant. We have made some comments on this situation in the main text. We have neglected $W_L - W_R$ mixing which is $\leq \mathcal{O}(10^{-3})$ and would cause a further suppression.

VII. ACKNOWLEDGMENT

S.G. would like to thank James Barry, Subrata Khan, Manimala Mitra and Werner Rodejohann for helpful discussions. S. Patra would like to thank the hospitality at PRL, where most of the present work was done.

-
- [1] S. Weinberg, Phys. Rev. Lett. **43**, 1566 (1979).
 - [2] H. V. Klapdor-Kleingrothaus [Heidelberg-Moscow and GENIUS Collaboration].
 - [3] W. Rodejohann, Int. J. Mod. Phys. E **20**, 1833 (2011).
 - [4] F. T. Avignone, III, S. R. Elliott and J. Engel, Rev. Mod. Phys. **80**, 481 (2008) [arXiv:0708.1033 [nucl-ex]]; J. J. Gomez-Cadenas, J. Martin-Albo, M. Mezzetto, F. Monrabal and M. Sorel, Riv. Nuovo Cim. **35**, 29 (2012) [arXiv:1109.5515 [hep-ex]].
 - [5] H. V. Klapdor-Kleingrothaus and I. V. Krivosheina, Mod. Phys. Lett. A **21**, 1547 (2006).
 - [6] H. V. Klapdor-Kleingrothaus, I. V. Krivosheina, A. Dietz and O. Chkvorets, Phys. Lett. B **586**, 198 (2004).
 - [7] M. Mitra, G. Senjanovic and F. Vissani, Nucl. Phys. B **856**, 26 (2012).
 - [8] A. Ibarra, E. Molinaro and S. T. Petcov, Phys. Rev. D **84**, 013005 (2011).
 - [9] S. Pascoli and S. T. Petcov, Phys. Rev. D **77**, 113003 (2008).
 - [10] M. Hirsch, H. V. Klapdor-Kleingrothaus and S. G. Kovalenko, Phys. Rev. D **53**, 1329 (1996)

- [hep-ph/9502385].
- [11] M. Hirsch, H. V. Klapdor-Kleingrothaus and S. G. Kovalenko, Phys. Rev. Lett. **75**, 17 (1995).
 - [12] A. Faessler, S. Kovalenko and F. Simkovic, Phys. Rev. D **58**, 055004 (1998) [hep-ph/9712535].
 - [13] H. Pas, M. Hirsch and H. V. Klapdor-Kleingrothaus, Phys. Lett. B **459**, 450 (1999) [hep-ph/9810382].
 - [14] G. Bhattacharyya, H. V. Klapdor-Kleingrothaus, H. Pas and A. Pilaftsis, Phys. Rev. D **67**, 113001 (2003).
 - [15] B. C. Allanach, C. H. Kom and H. Pas, Phys. Rev. Lett. **103**, 091801 (2009) [arXiv:0902.4697 [hep-ph]].
 - [16] V. Tello, M. Nemevsek, F. Nesti, G. Senjanovic and F. Vissani, Phys. Rev. Lett. **106**, 151801 (2011).
 - [17] M. Nemevsek, F. Nesti, G. Senjanovic and V. Tello, arXiv:1112.3061 [hep-ph].
 - [18] V. Cirigliano, A. Kurylov, M. J. Ramsey-Musolf and P. Vogel, Phys. Rev. Lett. **93**, 231802 (2004).
 - [19] W. Y. Keung and G. Senjanovic, Phys. Rev. Lett. **50**, 1427 (1983).
 - [20] G. Senjanovic and R. N. Mohapatra, Phys. Rev. D **12**, 1502 (1975). R. N. Mohapatra and J. C. Pati, Phys. Rev. D **11**, 2558 (1975). R. N. Mohapatra and J. C. Pati, Phys. Rev. D **11**, 566 (1975).
 - [21] P. Minkowski, Phys. Lett. B **67**, 421 (1977); T. Yanagida, proceedings of the *Workshop on Unified Theories and Baryon Number in the Universe*, Tsukuba, 1979, eds. A. Sawada, A. Sugamoto, KEK Report No. 79-18, Tsukuba; S. Glashow, in *Quarks and Leptons, Cargèse 1979*, eds. M. Lévy. et al., (Plenum, 1980, New York); M. Gell-Mann, P. Ramond, R. Slansky, proceedings of the *Supergravity Stony Brook Workshop*, New York, 1979, eds. P. Van Nieuwenhuizen, D. Freeman (North-Holland, Amsterdam); R. N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. **44**, 912 (1980).
 - [22] R. N. Mohapatra and G. Senjanovic, Phys. Rev. D **23**, 165 (1981); G. Lazarides, Q. Shafi and C. Wetterich, Nucl. Phys. B **181**, 287 (1981); J. Schechter and J. W. F. Valle, Phys. Rev. D **22**, 2227 (1980).
 - [23] R. Foot and H. Lew, Phys. Rev. D **42**, 945 (1990); R. Foot and G. Filewood, Phys. Rev. D **60**, 115002 (1999); M. Frank, A. Hayreter and I. Turan, Phys. Rev. D **83**, 035001 (2011); J. N. Esteves, J. C. Romao, M. Hirsch, W. Porod, F. Staub and A. Vicente, JHEP **1201**, 095

- (2012).
- [24] C. Y. Chen and P. S. B. Dev, arXiv:1112.6419 [hep-ph].
 - [25] J. Chakraborty, J. Gluza, R. Sevilano, R. Szafron' arXiv:1204.0736 [hep-ph].
 - [26] F. P. An *et. al.* [DAYA-BAY Collaboration], arXiv:1203.1669 [hep-ex].
 - [27] J. K. Ahn *et. al.* [Soo-Bong Kim for RENO collaboration], arXiv:1204.0626v1 [hep-ex]
 - [28] G. L. Fogli, E. Lisi, A. Marrone, A. Palazzo and A. M. Rotunno, Phys. Rev. D **84**, 053007 (2011) [arXiv:1106.6028 [hep-ph]].
 - [29] M. Hirsch, H. V. Klapdor-Kleingrothaus and O. Panella, Phys. Lett. B **374**, 7 (1996) [hep-ph/9602306].
 - [30] G. G. Senjanovic, Nucl. Phys. B **153**, 334 (1979).
 - [31] A. Pilaftsis, Z. Phys. C **55**, 275 (1992) [hep-ph/9901206]. J. Kersten and A. Y. Smirnov, Phys. Rev. D **76**, 073005 (2007) [arXiv:0705.3221 [hep-ph]]; R. Adhikari and A. Raychaudhuri, Phys. Rev. D **84**, 033002 (2011) [arXiv:1004.5111 [hep-ph]].
 - [32] S. T. Petcov, H. Sugiyama and Y. Takanishi, Phys. Rev. D **80**, 015005 (2009).
 - [33] W. Grimus and L. Lavoura, JHEP **0011**, 042 (2000) [hep-ph/0008179].
 - [34] Z. -z. Xing, Phys. Lett. B **679**, 255 (2009) [arXiv:0907.3014 [hep-ph]].